

Intersecting Branes, SUSY Breaking and the 2 TeV Excess at the LHC

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Abstract

Intersecting D-brane models in string theory can naturally support the gauge and matter content of left-right symmetric extensions of the Standard Model with gauge symmetry $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. Considering such models as candidates for explaining the 2 TeV excesses seen in Run-1 by both ATLAS and CMS, the minimal possible scale of supersymmetry breaking is determined by the requirement of precise one-loop gauge coupling unification. For the vector-like, bifundamental and (anti-)symmetric Higgs content of such brane configurations, this comes out fairly universally at around 19 TeV. For the $SU(2)_R$ gauge coupling one finds values $0.48 < g_R(M_R) < 0.6$. Threshold corrections can potentially lower the scale of supersymmetry breaking.

1 Introduction

In the 8 TeV run of the LHC the Higgs particle was discovered, which completes the particle spectrum of the Standard Model. Concerning physics beyond the SM, the hope was that supersymmetry would be found at a scale not far above the weak scale, thus providing a solution to the hierarchy problem. Even though no direct indication of supersymmetry has appeared yet, some anomalies were reported indicating with $2\text{--}3\sigma$ a resonance of around 2 TeV in the di-boson decay channel by ATLAS [1] and in the e^+e^-jj , Wh^0 and jj final states by CMS [2]. The most significant one is the ATLAS 3.4σ excess in the hadronic decay of a WZ pair of electro-weak gauge bosons.

Promising candidates to explain these excesses are left-right symmetric extensions of the Standard Model (LRSM) [3] with gauge symmetry $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ in which a bi-doublet Higgs field naturally generates mixings between the left and right $SU(2)$ W-bosons. The breaking of $SU(2)_R \times U(1)_{B-L}$ to the Standard Model hypercharge at the scale $M_R \sim 2\text{ TeV}$ can be performed e.g. by an $SU(2)_R$ triplet or simply by a doublet. The phenomenological potential of such models to explain all the excesses seen by ATLAS and CMS were e.g. analyzed in [4]. In these papers, the $SU(2)_R$ gauge coupling was restricted to be in the regime $0.4 < g_R(M_R) < 0.6$.

Taking the precision measurements for dilepton channels into account, some authors came to the conclusion that one should better use a leptophobic version of the LRSM, where the right handed leptons are not sitting in a doublet of $SU(2)_R$. Formally, this makes the model less natural so that our attitude is that both experimentally and phenomenologically this issue is not yet completely settled so that here we proceed to consider the standard version of an LRSM. Then, right-handed neutrinos and the Z' gauge boson should have masses larger than $M_{W'}$.

In fact, in this letter we are not so much concerned with such phenomenological fine-print, but would like to evaluate the potential of string theory inspired LRSM-like models to provide the overarching structure in the regime between M_R and a unification scale M_U . Let us mention that a string inspired explanation for the 2 TeV excesses was already proposed in terms of an anomalous Z' gauge boson in [5] (see also [6] for application to LHCb $b \rightarrow s\ell^+\ell^-$ anomalies). Note that there the mass of the anomalous Z' is of the order of the string scale M_S , which then is bound to be small.

In this letter we follow a different route and consider stringy realizations of LRSMs. One option is to embed the LRSM into a grand unified gauge group $SO(10)$, as it can be realized in heterotic or F-theory compactifications. In these constructions, one often gets a bunch of extra particles that can generate fast proton decay and one needs to solve the problem of mass hierarchies inside the Higgs multiplets (like the doublet-triplet splitting problem for $SU(5)$ GUTs).

Since all the fields in the LRSM sit in bifundamental and (anti-)symmetric

representations of the gauge group, it is also very natural to realize such models from intersecting D-branes [7] (see [8] for reviews), thereby avoiding some of the above mentioned problems present in GUT models. Baryon number is known to be a perturbative global symmetry in these constructions protecting the proton from a too fast decay. Therefore, in this paper we consider LRSM realizations by intersecting branes focusing on D7-branes in type IIB orientifolds.

Since in controllable models of string theory, supersymmetry needs to be broken at a scale M_{SUSY} smaller than the string scale $M_S \sim M_U$, we consider the hierarchy of mass scales $M_Z < M_R < M_{\text{SUSY}} < M_U$ and analyze the issue of gauge coupling unification. Such an analysis has been performed in a field theory context (see e.g. [9, 10] for more recent studies), but to our knowledge it was always assumed either that one does not have supersymmetry below the GUT scale at all or that the intermediate hierarchy of scales is reversed, i.e. $M_Z < M_{\text{SUSY}} < M_R < M_U$. Thus, the observations made in this letter can be considered as complementary to earlier results reported in the literature.

This letter is organized as follows: In section 2, we review a few important aspects of intersecting D-brane models where in particular we discuss under what circumstance one can still get gauge coupling unification at the string scale. Section 3 provides two simple prototype examples of possible realizations of a LRSM.

In section 4 the one-loop running of the four gauge couplings is analyzed. Note that this analysis only depends on the matter content of the models and is therefore generically valid, i.e. without necessarily referring to an intersecting D-brane scenario. After fixing the Higgs sector, requiring precise one-loop gauge coupling unification and using $M_R = 2 \text{ TeV}$, one can uniquely determine the unification scale, the supersymmetry breaking scale and the $SU(2)_R$ gauge coupling g_R . It is observed that the minimal possible value for the supersymmetry breaking scale, 19 TeV, shows a certain universal behavior. Moreover, the value of the $SU(2)_R$ gauge coupling come out as $0.48 < g_R(M_R) < 0.6$, depending on the vector-like Higgs sector.

Note added: The latest announcement of physics results by the ATLAS+CMS collaborations on 15.12.2015 did not show any 2 TeV excesses at Run-2. This of course diminishes the experimental motivation for the analysis done in this letter, but we think that the general results on the relation of the scales in LRSMs and supersymmetry that are obtained are nevertheless interesting.

2 Models of intersecting D7-branes

Let us review some of the main ingredients for the construction of intersecting D7-brane models in orientifolds of the Type IIB superstring. For more details we refer to the reviews [8] and in particular to [11].

One considers the Type IIB superstring compactified on a Calabi-Yau three-fold X and performs the quotient by an orientifold projection $\Omega\sigma(-1)^{F_L}$, where

Ω is the world-sheet parity reversal and σ denotes a holomorphic involution of the Calabi-Yau satisfying $\sigma(J) = J$ and $\sigma(\Omega_3) = -\Omega_3$. Here J is the Kähler form and Ω_3 is the holomorphic $(3,0)$ form on the threefold. At the fixed locus of σ , one gets O7 and O3-planes, whose R-R charges (tadpoles) need to be cancelled by the introduction of D7 and D3-branes.

Stacks of N_a D7-branes can wrap holomorphic four-cycles Σ_a in the homology class $H_4(X, \mathbb{Z})$ of the threefold. Moreover, these D7-branes can carry a non-trivial gauge flux determining a line bundle \mathcal{L}_a on Σ_a . Depending on whether this D7-brane configuration is invariant under the orientifold projection or not, one gets $SO(N)/SP(N)$ -gauge groups or $U(N)$ gauge groups, respectively. Note that one gets $U(N) = SU(N) \times U(1)$ instead of $SU(N)$.

The most simple and phenomenologically useful configurations for SM-like model building arise for branes wrapping orientifold invariant, rigid four-cycles Σ_a of SP -type with the different stacks being only distinguished by their gauge flux \mathcal{L}_a . If $c_1(\mathcal{L}_a) \in H_-^2(\Sigma_a, \mathbb{Z})$ then such a stack carries gauge group $SP(N_a)$ and for a flux $c_1(\mathcal{L}_a) \in H_+^2(\Sigma_a, \mathbb{Z})$ one gets an orientifold image line bundle \mathcal{L}'_a so that together they support a $U(N_a)$ gauge group.

The massless modes between two such stacks of branes transform in bifundamental or (anti)-symmetric representations, where the chiral index is given by

$$I_{ab} = - \int_X [\Sigma_a] \wedge [\Sigma_b] \wedge (c_1(L_a) - c_1(L_b)) . \quad (2.1)$$

Here $[\Sigma]$ denotes the Poincaré dual two form to the four-cycle Σ . For $U(N)$ stacks of branes the various open string sectors give the chiral spectrum summarized in table 1.

sector	$U(N_a)$	$U(N_b)$	chirality
(ab)	$\bar{\square}_{(-1)}$	$\square_{(1)}$	I_{ab}
$(a'b)$	$\square_{(1)}$	$\square_{(1)}$	$I_{a'b}$
$(a'a)$	$\square_{(2)}$	1	$\frac{1}{2}(I_{a'a} + 2I_{O7a})$
$(a'a)$	$\square\square_{(2)}$	1	$\frac{1}{2}(I_{a'a} - 2I_{O7a})$

Table 1: Chiral spectrum for intersecting D7-branes. The subscripts denote $U(1)$ charges.

For an SP -stack, chiral fields can only arise from the intersection with a $U(N)$ stack.

The non-chiral (vector-like) part of the spectrum can be determined by computing certain line bundle cohomology groups on Σ ¹. For branes wrapping the

¹For concrete model building it is important to take into account that the four-cycle Σ

same four-cycle Σ , this was determined in [11] and for completeness we provide the result

$$\begin{aligned}\text{Ext}^0(1_*L_a, 1_*L_b) &= H^0(\Sigma, L_a \otimes L_b^\vee), \\ \text{Ext}^1(1_*L_a, 1_*L_b) &= H^1(\Sigma, L_a \otimes L_b^\vee) + H^0(\Sigma, L_a \otimes L_b^\vee \otimes N_D), \\ \text{Ext}^2(1_*L_a, 1_*L_b) &= H^2(\Sigma, L_a \otimes L_b^\vee) + H^1(\Sigma, L_a \otimes L_b^\vee \otimes N_D), \\ \text{Ext}^3(1_*L_a, 1_*L_b) &= H^2(\Sigma, L_a \otimes L_b^\vee \otimes N_D).\end{aligned}\tag{2.2}$$

For branes wrapping two different 4-cycles intersecting over a curve $C = \Sigma_a \cap \Sigma_b$, one has to compute the cohomology classes

$$H^i(C, L_a^\vee \otimes L_b \otimes K_C^{\frac{1}{2}}), \quad i = 0, 1.\tag{2.3}$$

For more details, we refer to [11].

Moreover, the $U(1)$ factors in the total gauge group are often anomalous. In string theory, these anomalies are cancelled by the Green-Schwarz mechanism, where a shift symmetry of an axion gets gauged so that it becomes the longitudinal mode of a massive abelian vector field. The abelian gauge symmetry gets broken but survives as a perturbative global symmetry that is only broken by D-brane instantons. We recall that an abelian gauge field can even gain a mass without having an anomaly.

Let us consider now the tree-level gauge couplings at the string scale

$$\kappa_a \frac{4\pi}{g_a^2} = \tau_a - \frac{1}{2g_s} \int_{\Sigma} c_1^2(\mathcal{L}_a)\tag{2.4}$$

where $\tau_a = \frac{1}{2\ell_s^4} \int_{\Sigma_a} J \wedge J$ denotes the volume of the four-cycle (in Einstein frame) in units of the string length ℓ_s and $g_s = e^\varphi$ is the string coupling constant. Here $\kappa_a = 1$ for a $U(N)$ stack and $\kappa_a = 2$ for an $SP(2N)/SO(2N)$ stack². As usual in string theory, these couplings will receive one-loop threshold corrections. Note that for all branes wrapping the same 4-cycle, the differences among the gauge couplings g_a come only from the line bundles and in fact only from the value of $\int_{\Sigma_a} c_1^2(\mathcal{L}_a) \in \mathbb{Z}$ ³.

In contrast to single GUT groups, the gauge couplings at the string scale are not necessarily equal. This situation and its consequences for gauge coupling unification for intersecting D-brane models and F-theory models have been discussed in [12]. However, one can still design situations where precise unification can occur.

can contain homological two-cycles that are not closed in the ambient threefold X . One can always twist a line bundle \mathcal{L}_a by such a trivial bundle \mathcal{R}_a without changing the chiral spectrum. However, the vector-like part computed via (2.2) for the bundle $\mathcal{L}_a \otimes \mathcal{R}_a$ will change.

²The factor $\kappa = 1, 2$ is due to the “doubled” embedding of $U(N)$ into $SO(2N)/SP(2N)$

$T_{2N} = \frac{1}{\sqrt{2}} \begin{pmatrix} t_N & 0 \\ 0 & -t_N^T \end{pmatrix}.$

³ These values also depend on the twist line bundle \mathcal{R}_a .

If all branes are of $U(N)$ type, one can wrap all stacks around a single 4-cycle Σ of size τ and distinguish them solely by the line-bundles \mathcal{L}_a .

- In the regime $\tau \gg g_s^{-1} > 1$ with the differences of the flux integrals not being too large, the gauge couplings approximately unify with $\frac{4\pi}{g_a^2} = \tau$ where the flux dependent corrections can be considered as small threshold corrections. The values of τ and g_s are determined by moduli stabilization.
- Since the second term in (2.4) only depends on the topological quantity $\text{ch}_2(\mathcal{L}_a)$, all gauge couplings can still be degenerate even if the \mathcal{L}_a themselves are different.

If some of the branes are of SP -type then due to the factor $\kappa = 2$ in (2.4), the SP branes should wrap a different four-cycle Σ_{SP} than the $U(N)$ branes Σ_U with $\tau_{SP} = 2\tau_U$. In this case, one also gets approximate gauge coupling unification in the sense just explained for $U(N)$ stacks.

Thus, in the following we will work in a scheme where we realize the LRSM on such intersecting D7-branes and we will assume that we have gauge coupling unification at the string or unification scale up to small threshold corrections. Moreover, the initial brane realization should be supersymmetric at the string scale, where, as usual, supersymmetry breaking will be mediated to the observable sector by generating soft masses of the order M_{SUSY} .

3 Brane realizations of LRSM

In this section two principal realizations of the LRSM in terms of intersecting D7-branes are presented. We do not provide fully fledged global string compactifications, but instead only local brane configurations that satisfy the consistency conditions following from string theory. The realization of the SM itself in terms of intersecting branes is also known as the Madrid quiver, first presented in [13]. That model is very similar to the LRSM quivers [14] to be discussed below.

LRSM quiver A

First we consider the simplest quiver A, in which the $SU(2)_L \times SU(2)_R$ sector is realized on orientifold invariant branes supporting $SP(2) \simeq SU(2)$ gauge group. We introduce four stacks of D7-branes carrying appropriate line bundles such that one gets the initial gauge symmetry $U(3) \times SP(2) \times SP(2) \times U(1)$ and that the massless spectrum is the one shown in table 2.

One has two abelian gauge factors $U(1)_B = \frac{1}{3}U(1)_a \subset U(3)$ and $U(1)_L = U(1)_d$, whose charges can be identified with baryon and lepton number, respectively. However, the only anomaly-free combination is $U(1)_{B-L} = \frac{1}{3}U(1)_a - U(1)_d$. Therefore, the orthogonal combination receives a mass via the Green-Schwarz

number	field	$SU(3)$	$SU(2)_L$	$SU(2)_R$	$U(1)_a \times U(1)_d$	$U(1)_{B-L}$
3	Q_L	3	2	1	(1, 0)	1/3
3	$(Q_R)^c$	$\bar{3}$	1	2	(-1, 0)	-1/3
3	ℓ_L	1	2	1	(0, 1)	-1
3	$(\ell_R)^c$	1	1	2	(0, -1)	1
N_Φ	Φ	1	2	2	(0, 0)	0
N_H	H_R^u	1	1	2	(0, 1)	-1
N_H	H_R^d	1	1	2	(0, -1)	1
N_Δ	Δ_0	1	1	3	(0, 0)	0
N_S	S^u	1	1	1	(0, 2)	-2
N_S	S^d	1	1	1	(0, -2)	2

Table 2: Massless left-handed spectrum of LRSM quiver A.

mechanism. This in particular means that baryon and lepton number survive as global symmetries and can protect the proton from decaying. In particular, the unification scale can be smaller than the usual GUT-scale $M_{\text{GUT}} = 2 \cdot 10^{16}$ GeV.

The gauge coupling of a linear combination $U(1)_{B-L} = \sum c_i U(1)_i$ with $U(1)_i \subset U(N_i)$ can be computed as

$$\frac{1}{\alpha_{B-L}} = \sum_i \frac{1}{2} N_i c_i^2 \frac{1}{\alpha_i} = \frac{1}{2} \left(\frac{1}{3} \frac{1}{\alpha_a} + \frac{1}{\alpha_d} \right) = \frac{2}{3} \frac{1}{\alpha_s}, \quad (3.1)$$

assuming stringy gauge coupling unification, i.e. $\alpha_s = \alpha_a = \alpha_d$. Note that this is the same relation as for $SO(10)$ GUTs.

The bi-doublet Higgs field Φ originates from a vector-like intersection between the two $SP(2)$ -branes. As indicated in the table, there could be more than just a single such Higgs field, but its minimal non-vanishing number is really $N_\Phi = 1$. It is clear that for intersecting branes, one cannot get an $SU(2)_R$ triplet with $Q_{B-L} = \pm 2$. The open string of such a massless mode would need four instead of two ends. Therefore, the breaking of the $SU(2)_R \times U(1)_{B-L}$ gauge symmetry to $U(1)_Y$ has to proceed via a Higgs field in the doublet representation of $SU(2)_R$ with $Q_{B-L} = 1$. Note that anomaly cancellation forces us here to introduce such Higgs fields in vector-like pairs $H_R^{u,d}$. We also added vector-like pairs $S^{u,d}$ of fields transforming in the symmetric representation of $U(1)_d$ and $SU(2)_R$ triplets with $Q_{B-L} = 0$. Note that in contrast to other approaches, parity symmetry $P : SU(2)_L \leftrightarrow SU(2)_R$ is broken explicitly in this Higgs sector.

LRSM quiver B

One can also realize the $SU(2)_{L,R}$ gauge symmetries on $U(2)$ type of branes. In this case one has four $U(1)$ factors, of which we assume that only the anomaly-free combination $U(1)_{B-L} = \frac{1}{3}U(1)_a - U(1)_d$ stays massless after the Green-Schwarz mechanism has been employed. A configuration consistent with the

stringy constraints is presented in table 3. Note that the generation of all possible SM Yukawa couplings and anomaly cancellation forces one to introduce an even number of bi-doublet Higgses Φ .

number	field	$SU(3)$	$SU(2)_L$	$SU(2)_R$	$U(1)^4$	$U(1)_{B-L}$
2	Q_L	3	2	1	(1, 1, 0, 0)	1/3
1	Q_L	3	2	1	(1, -1, 0, 0)	1/3
2	$(Q_R)^c$	$\bar{3}$	1	2	(-1, 0, 1, 0)	-1/3
1	$(Q_R)^c$	$\bar{3}$	1	2	(-1, 0, -1, 0)	-1/3
3	ℓ_L	1	2	1	(0, -1, 0, 1)	-1
3	$(\ell_R)^c$	1	1	2	(0, 0, -1, -1)	1
$N_\Phi/2$	Φ^u	1	2	2	(0, 1, 1, 0)	0
$N_\Phi/2$	Φ^d	1	2	2	(0, -1, -1, 0)	0
N_H	H_R^u	1	1	2	(0, 0, 1, 1)	-1
N_H	H_R^d	1	1	2	(0, 0, -1, -1)	1
$N_\Delta/2$	Δ_0^u	1	1	3	(0, 0, 2, 0)	0
$N_\Delta/2$	Δ_0^d	1	1	3	(0, 0, -2, 0)	0
N_S	S^u	1	1	1	(0, 0, 0, 2)	-2
N_S	S^d	1	1	1	(0, 0, 0, -2)	2

Table 3: Massless left-handed spectrum of LRSM quiver B. Anomaly cancellation/tadpole cancellation enforces N_Φ even.

In the next section, we analyze the running of the four gauge couplings in the regime $M_R < M_{\text{SUSY}} < M_U$. Even though the matter content of the models was motivated by D-brane constructions, the upcoming analysis only depends on the former and is therefore generically valid.

4 Gauge Coupling Unification

We first run the Standard Model couplings from the weak scale up to the new left-right unification scale $M_R \sim 2 \text{ TeV}$. For the values of the gauge couplings at the weak scale $M_Z = 91.18 \text{ GeV}$ we took $\alpha_s = 0.1172$, $\alpha = 1/127.934$ and $\sin \theta_w = 0.23113$. Then, at the scale M_R one obtains

$$\alpha_s(M_R) = 0.0835, \quad \alpha_L(M_R) = 0.0321, \quad \alpha_Y(M_R) = 0.0105. \quad (4.1)$$

At the scale M_R the hypercharge coupling splits into the $SU(2)_R$ and the $U(1)_{B-L}$ coupling according to

$$\frac{1}{\alpha_Y} = \frac{1}{\alpha_R} + \frac{1}{\alpha_{B-L}}. \quad (4.2)$$

The running beyond M_R is evaluated under the following two assumptions

1. Following the extended survival hypothesis [15], in the regime $M_R < \mu < M_{\text{SUSY}}$ there is just the minimal particle content of the non-supersymmetric LRSM, i.e. in particular one scalar Higgs bi-doublet Φ and one scalar Higgs doublet H_R^u . Due to supersymmetry breaking, they are expected to gain soft masses of the order of M_{SUSY} . Since for $M_{\text{SUSY}} \gg M_R$ one cannot refer to supersymmetry to solve the hierarchy problem, one needs some fine-tuning or string landscape argument to achieve this.
2. In the regime $M_{\text{SUSY}} < \mu < M_U$ all the supersymmetric states that the intersecting D-brane model provides contribute to the running. This includes N_Φ chiral fields in the bi-doublet representation, as well as N_H vector-like Higgs fields H_R . We also allow for N_Δ vector-like fields Δ_0 and N_S vector-like fields S .

The one-loop running of the four gauge couplings occurs according to

$$\frac{1}{\alpha_i(\mu)} = \frac{1}{\alpha_i(M_R)} + \frac{b_i}{2\pi} \log \left(\frac{M_{\text{SUSY}}}{M_R} \right) + \frac{\tilde{b}_i}{2\pi} \log \left(\frac{\mu}{M_{\text{SUSY}}} \right) \quad (4.3)$$

where $(\alpha_3, \alpha_2^L, \alpha_2^R, \alpha_1) = (\alpha_s, \alpha_L, \alpha_R, \frac{2}{3}\alpha_{B-L})$ and the β -function coefficients in the non-supersymmetric and supersymmetric region are given by⁴

$$\begin{aligned} b_3 &= 7, & \tilde{b}_3 &= 3, \\ b_2^L &= 3, & \tilde{b}_2^L &= -N_\Phi, \\ b_2^R &= -\frac{1}{6}n_H^l + 3, & \tilde{b}_2^R &= -N_H - N_\Phi - 2N_\Delta, \\ b_1 &= -\frac{1}{4}n_H^l - 4, & \tilde{b}_1 &= -\frac{3}{2}N_H - 3N_S - 6. \end{aligned} \quad (4.4)$$

Here we were leaving the number n_H^l of light scalar Higgs fields in the doublet representation of $SU(2)_R$ after supersymmetry breaking as an open parameter. A fermion from the superfield H_R that remains light contributes like $n_H^l = 2$. The extended survival hypothesis means that we have $n_H^l = 1$.

Taking the relation (4.2) into account, we have three unspecified parameters namely $\alpha_R(M_R)$, M_{SUSY} and M_U that can be uniquely determined by requiring that all four gauge couplings unify at a scale M_U . For self-consistency one needs $M_R < M_{\text{SUSY}} < M_U$ and that all couplings remain in the perturbative regime. Let us discuss the two classes $N_\Phi = \text{even/odd}$ separately.

LRSM quiver A for odd N_Φ

For $N_\Delta = N_S = 0$, choosing different values of the number of Higgs fields N_Φ and N_H and solving for g_R , M_{SUSY} and M_U , for quiver A one obtains the scales shown in table 4.

⁴Due to the normalization in (3.1), one has e.g. $\tilde{b}_1 = \frac{3}{2} \sum_{i, \text{chiral}} \frac{Q_{B-L,i}^2}{4}$.

N_H	$M_{\text{SUSY}} [\text{GeV}]$	g_R	N_H	$M_{\text{SUSY}} [\text{GeV}]$	g_R	$M_U [\text{GeV}]$
< 3	$< M_R$		< 6	$< M_R$		
3	$19.2 \cdot 10^3$	0.532	6	$19.2 \cdot 10^3$	0.507	$2.18 \cdot 10^{12}$
4	$1.26 \cdot 10^8$	0.537	7	$1.34 \cdot 10^7$	0.516	$1.94 \cdot 10^{13}$
5	$1.21 \cdot 10^{10}$	0.540	10	$7.70 \cdot 10^8$	0.522	$7.47 \cdot 10^{13}$
6	$1.99 \cdot 10^{11}$	0.542	12	$1.21 \cdot 10^{10}$	0.526	$1.87 \cdot 10^{14}$
\downarrow			\downarrow			
∞	M_{GUT}	~ 0.55	∞	M_{GUT}	~ 0.55	M_{GUT}

Table 4: M_{SUSY} , g_R and M_U for different values of N_H with $N_\Delta = N_S = 0$. The left table is for $N_\Phi = 1$ ($M_U = M_{\text{GUT}}$) and the right one for $N_\Phi = 3$.

For $N_H \leq 3, 6$ one gets $M_{\text{SUSY}} < M_R$ which is in conflict with the assumption. The lowest possible values are therefore $N_H = 3, 6$, for which the supersymmetry breaking scale comes out one order of magnitude larger than M_R and is actually the same for the two choices $N_\Phi = 1$ and $N_\Phi = 3$.

This value would be out of reach of the LHC Run-2, but threshold corrections at the high scale and two-loop corrections to the running might lower this value. This issue will be discussed below. The value of the $SU(2)_R$ gauge coupling does only vary slightly in the region $0.5 < g_R < 0.55$. Recall that baryon number is still a perturbative global symmetry so that for $M_U < M_{\text{GUT}}$ fast proton decay is not an issue. In figure 1, the running coupling constants are shown for $N_\Phi = 1$ and $N_H = 3$.

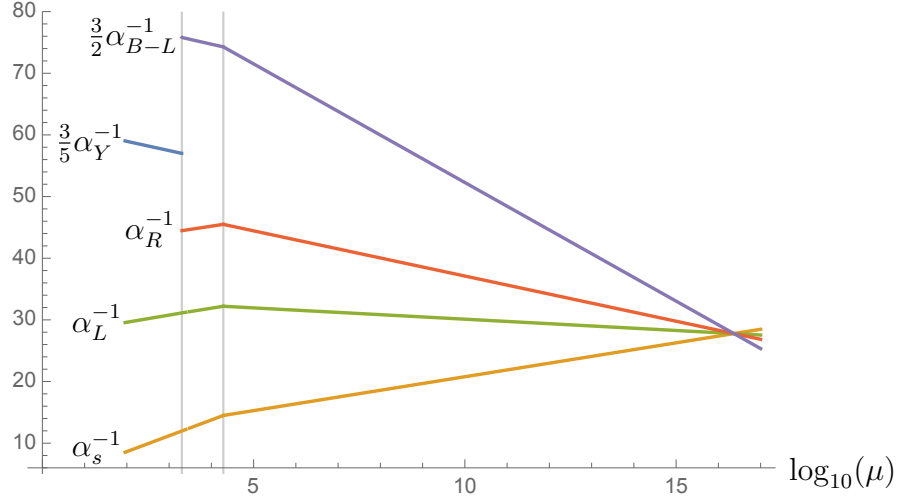


Figure 1: One-loop running of the LRS gauge couplings for $N_\Phi = 1$, $N_H = 3$, $N_\Delta = N_S = 0$ and $n_H^l = 1$ with $M_{\text{SUSY}} = 19.2 \text{ TeV}$ and $g_R = 0.532$.

We observed that the minimal value of the supersymmetry breaking scale $M_{\text{SUSY}} = 19.2 \text{ TeV}$ was appearing for both choices of the number of bi-doublet

Higgses $N_\Phi = 1, 3$. Clearly, this asks for an explanation. First, one notices that the values of M_{SUSY} and M_U do only depend on the combination $(N_H + N_\Delta + N_S)$ with only g_R depending on their individual values. For the (minimal) choice $(N_H + N_\Delta + N_S)^{\min} = \frac{3}{2}(N_\Phi + 1)$, the supersymmetry breaking scale can be generically determined as

$$M_{\text{SUSY}} = M_R \exp \left[\frac{2\pi}{14 - n_H^l} \left(\frac{12}{\alpha_L(M_R)} - \frac{7}{\alpha_s(M_R)} - \frac{3}{\alpha_Y(M_R)} \right) \right]. \quad (4.5)$$

Surprisingly, the scale does not depend on N_Φ , but only on the LR-breaking scale M_R and the number of light Higgses. For $n_H^l = 1$ one gets $M_{\text{SUSY}} = 19.2 \text{ TeV}$, which increases for larger values of n_H^l . For $M_R = 1.8 \text{ TeV}$ one finds $M_{\text{SUSY}} = 16 \text{ TeV}$.

One can show that for $N_\Phi \leq 21$, the number $(N_H + N_\Delta + N_S)^{\min} = \frac{3}{2}(N_\Phi + 1)$ is indeed the minimal threshold value guaranteeing $M_{\text{SUSY}} > M_R$. It is in this sense that, for $M_R = 2 \text{ TeV}$, this is a universal result. For $N_\Phi = 21$, $N_\Delta = N_S = 0$ and $n_H^l = 1$ one obtains $M_{\text{SUSY}} = 19.2 \text{ TeV}$, $g_R = 0.476$ and $M_U = 1.98 \cdot 10^6 \text{ GeV}$, hence featuring smaller values of the $SU(2)_R$ gauge coupling and the unification scale.

Taking the minimal choice $N_\Phi = 1$, for all possible partitions $(N_H + N_\Delta + N_S)^{\min} = 3$ with $N_H \geq 1$ we find $M_{\text{SUSY}} = 19.2 \text{ TeV}$, $M_U = M_{\text{GUT}} = 2.3 \cdot 10^{16} \text{ GeV}$ and the values of $0.48 < g_R < 0.6$ shown in table 5.

(N_H, N_Δ, N_S)	g_R
(3, 0, 0)	0.532
(2, 1, 0)	0.507
(2, 0, 1)	0.560
(1, 2, 0)	0.485
(1, 1, 1)	0.532
(1, 0, 2)	0.594

Table 5: Values of g_R for $N_\Phi = 1$ and $(N_H + N_\Delta + N_S)^{\min} = 3$.

LRSM quivers A,B for even N_Φ

The same computation can be performed for an even number of bi-doublet Higgs fields. Recall that for quiver B, anomaly cancellation enforced these fields to come in vector-like pairs. In table 6 we list the resulting mass scales for various choices of N_Φ for the corresponding minimal value of N_H and $N_\Delta = N_S = 0$.

Here, we do not find the same universality as for N_Φ odd. This would arise for half-integer values of the number of vector-like Higgs fields H_R . For $2 \leq N_\Phi \leq 6$, the minimal choices of N_H are parameterized as $N_H^{\min} = \frac{3}{2}N_\Phi + 2$. For $N_\Phi = 8$ the branch changes to $N_H^{\min} = \frac{3}{2}N_\Phi + 1$.

N_Φ	N_H^{\min}	M_{SUSY} [GeV]	g_R	M_U [GeV]
2	5	$1.48 \cdot 10^6$	0.521	$2.12 \cdot 10^{14}$ GeV
4	8	$4.91 \cdot 10^5$	0.506	$6.19 \cdot 10^{11}$ GeV
6	11	$2.56 \cdot 10^5$	0.497	$1.90 \cdot 10^{10}$ GeV
8	13	$1.5 \cdot 10^3$	0.483	$9.40 \cdot 10^7$ GeV

Table 6: Values of the M_{SUSY} , g_R and M_U for different values of N_Φ and the corresponding minimal values of N_H with $n_H^l = 1$ and $N_\Delta = N_S = 0$.

What one can do though is to add N_T vector-like pairs of massless modes in the representation $(3, 1, 1, \pm 2/3)$. These arise from open strings stretched between the $U(3)$ stack and the $U(1)_d$ stack. The β -function coefficients in the susy regime $M_R < \mu < M_U$ then read

$$\begin{aligned}
\tilde{b}_3 &= 3 - N_T, \\
\tilde{b}_2^L &= -N_\Phi, \\
\tilde{b}_2^R &= -N_H - N_\Phi - 2N_\Delta, \\
\tilde{b}_1 &= -\frac{3}{2}N_H - 3N_S - 6 - N_T.
\end{aligned} \tag{4.6}$$

Then it is clear that choosing N_Φ even and $N_T = 1$, the three values M_{SUSY} , $g_R(M_R)$ and M_U remain the same as for $N_\Phi - 1$ bi-doublet Higgses and $N_T = 0$. Only the value of the unified gauge coupling changes. Hence, one is back to the discussion for odd N_Φ . In figure 2, the one-loop running coupling constants are shown for $N_\Phi = 2$, $N_T = 1$ and $N_H = 3$.

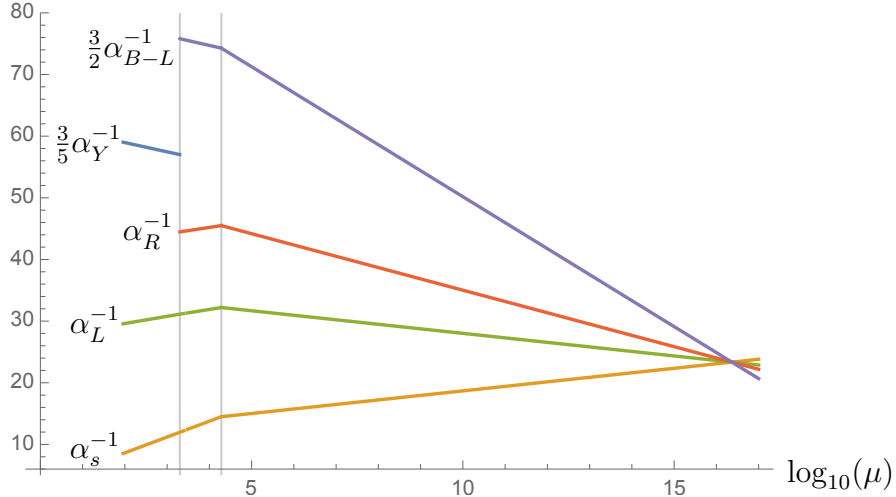


Figure 2: Running for $N_\Phi = 2$, $N_T = 1$, $N_\Delta = N_S = 0$ and $N_H = 3$ with $M_{\text{SUSY}} = 19.2$ TeV and $g_R = 0.532$.

Comment

The value of $M_{\text{SUSY}} = 19 \text{ TeV}$ is a bit too large to be detected directly at the LHC. However, our analysis was very strict in the sense that we were assuming that all masses at the scales M_R and M_{LRMSUSY} are the same, respectively. Moreover, the computation is performed only at one-loop level, where two-loop corrections are usually expected to give a 4% correction to the couplings at the unification scale. A correction of the same order is expected if the gauge fluxes on the stacks of branes lead to string threshold corrections, as discussed in section 2.

Just to get a first impression, for quiver A let us assume that e.g. the supersymmetry breaking scale is 3 TeV, with $g_R = 0.53$ and just run the couplings up to the GUT scale. The result is shown in figure 3.

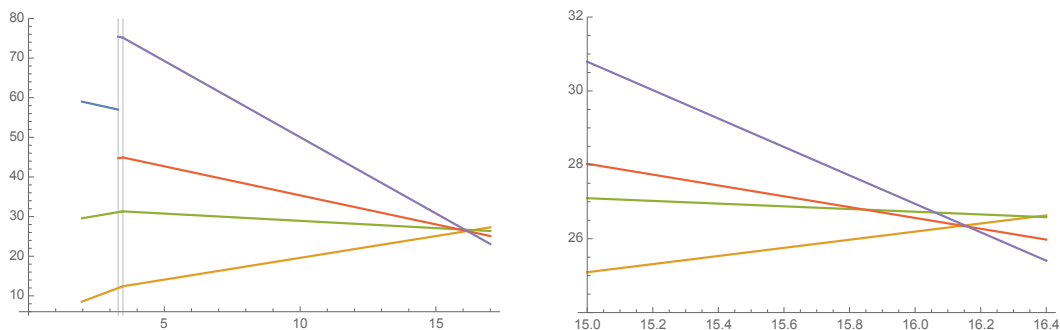


Figure 3: One-loop running of the LRS gauge couplings for $N_H = 3$ with $M_{\text{SUSY}} = 3 \text{ TeV}$ and $g_R = 0.53$.

In the left picture one does not even see that the couplings do not unify exactly. From the right picture one can estimate a 2% failure in doing so. Therefore, we conclude that for the minimal value $N_H = 3$, the inclusion of threshold corrections and probably also two-loop corrections can make a supersymmetry breaking scale of just a bit above M_R still consistent with stringy unification at around the GUT scale.

5 Conclusions

We presented possible realizations of LRSs in terms of intersecting D7-brane quivers. Since only bifundamental and (anti-)symmetric representations can occur, the SM and LRS Higgs representations were fairly constrained. Employing this point, we were considering the minimal matter content and were studying the running coupling constants in the regime between the LR-breaking scale of 2 TeV, as suggested by excesses in the LHC data, and a potential unification scale.

We found that for a not too large number of bi-doublet Higgses H_Φ and in each case a sufficiently large number of vector-like fields $\{H_R, \Delta_0, S\}$ one indeed achieves a precise one-loop unification scenario in which the minimal value of the supersymmetry breaking scale was determined as 19 TeV. This value was shown to be universal for an odd number of bi-doublet Higgs fields. For an even number of such Higgses, by adding a further vector-like colour triplet state, the same universality could be achieved⁵. Threshold corrections at the high scale might lower this value so that it can be detectable at the LHC Run-2.

Even though one can contemplate various variations and extensions of such stringy LRSM models, it is satisfying that with some stringy input and a number of reasonable assumptions it was possible to derive such a universal and in this scheme predictive result. Moreover, the value of the gauge coupling constant g_R was not varying much, either, and came out as $0.48 < g_R < 0.6$. This is in the regime that was also suggested by a more phenomenological analysis of the LRSM to fit the various, of course still not significant, 2 TeV excesses observed in Run-1 at the LHC.

It would be interesting to generalize the computation in various directions like e.g. to D-brane realizations of leptophobic models. One should also include 2-loop effects and consider threshold effects both at the small and the large scale.

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⁵We were informed that the model with $N_\Phi = 2$, $N_T = 1$, $N_H = 1$ and $N_S = 2$ has also been considered in [10].

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